

SUPERSYMMETRIC CP VIOLATION ε'/ε DUE TO ASYMMETRIC A-MATRIX

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ABSTRACT

We study contributions of supersymmetric CP phases to the CP violation ε'/ε in models with asymmetric A -matrices. We consider asymmetric A -matrices, which are obtained from string-inspired supergravity. We show that a certain type of asymmetry of A -matrices enhances supersymmetric contributions to the CP violation ε'/ε and the supersymmetric contribution to ε'/ε can be of order of the KTeV result, $\varepsilon'/\varepsilon \sim 10^{-3}$.

1. CP violation is sensitive to physics beyond the standard model (SM), that is, CP violation, as well as flavor changing neutral current (FCNC) processes, constrains significantly physics beyond the SM, *e.g.*, supersymmetric models. On the other hand, if any deviation of CP violation from the SM is observed, that would be a strong hint to physics beyond the SM. Many works have been done on CP violation in supersymmetric models [1, 2, 3, 4].

Recently the KTeV collaboration at Fermilab has reported a measurement of the direct CP-violation in $K \rightarrow \pi\pi$ decays [5]

$$\text{Re}(\varepsilon'/\varepsilon) = (28 \pm 4) \times 10^{-4}.$$

This measurement has confirmed the previous result of the NA31 experiment at CERN [6]. Hence, it excludes the superweak models. The SM predicts non-zero value for ε'/ε . However, its prediction suffers from a large ambiguity due to the theoretical uncertainties in the hadronic quantities. Recent discussions on $(\varepsilon'/\varepsilon)$ can be found in Refs.[7]-[12].

In supersymmetric models, there are several possible sources for CP violation beyond the Cabibbo-Kobayashi-Maskawa (CKM) phase δ_{CKM} in the SM. Two types of physical phases only remain in the minimal supersymmetric standard model (MSSM) after all appropriate field redefinition, namely, the phases of A -parameters and the phase of the μ -term (ϕ_μ). A generic type of A -matrices include many degree of freedom for the real part and imaginary part. However, in most of cases the universality of A -matrices has been assumed. That is good to simplify calculations, but that removes interesting degree of freedom, in particular for CP violation.

Several implications of non-universal A -matrices have been discussed in Ref. [13, 14]. The non-universality among the soft supersymmetry (SUSY) breaking terms plays an important role on all CP violating processes. In particular, it has been shown in Ref. [14] that non-degenerate A -parameters can generate the experimentally observed CP violation ε even with the vanishing CKM CP phase $\delta_{CKM} = 0$, that is, the fully supersymmetric CP violation in the Kaon system is possible. In Ref. [14], an example of symmetric A -matrices for the first 2×2 block has been considered with the exact degeneracy of squark masses between the first and second families. Because string-derived soft terms [15, 16] require symmetric A -matrices for exactly degenerate soft masses. This model can lead to the parameter region where we have the SUSY contribution of $O(10^{-3})$ to ε . However, this type of models provide with a very small value for ε'/ε . This result is due to an accidental cancellation between the different contributions because of the symmetric form of the A -matrices which have been used.

In this letter, we study the SUSY contributions to the CP violation ε'/ε in models with asymmetric A -matrices. We consider two possibilities for asymmetric A -matrices keeping the degeneracy of squark masses. One model has almost degenerate squark masses, which are realized by dilaton-dominate SUSY breaking. In the other model, we require a delicate cancellation between string-derived soft masses and the D -term contributions to soft masses. The latter can lead to a large asymmetry for the A -matrix. Using these models, we calculate ε'/ε explicitly. Then we show that in the case with asymmetric A -terms the

SUSY contribution to ε'/ε can be of order of the KTeV result, $\varepsilon'/\varepsilon \sim 10^{-3}$. In the whole analysis we take $\delta_{CKM} = 0$ in order to show the pure SUSY contributions.

2. Here we consider the possibilities that one can obtain non-degenerate A -matrices keeping degeneracy of the squark masses. First we give a brief review on the soft SUSY breaking terms in string models,

$$-\mathcal{L}_{\text{SB}} = \frac{1}{6} h_{ijk} \phi_i \phi_j \phi_k + \frac{1}{2} (\mu B)^{ij} \phi_i \phi_j + \frac{1}{2} (m^2)_i^j \phi^{*i} \phi_j + \frac{1}{2} M_a \lambda \lambda + \text{H.c.} \quad (1)$$

where the ϕ_i are the scalar parts of the chiral superfields Φ_i and λ are the gauginos. We use the notation of trilinear coupling terms, the so-called A -terms, as $h_{ijk} = (YA)_{ijk}$, where Y_{ijk} is the corresponding Yukawa coupling.

We start with the (weakly coupled) string-inspired supergravity theory. Its Kähler potential is

$$K = -\ln(S + S^*) + 3\ln(T + T^*) + \sum_i (T + T^*)^{n_i} |\Phi^i|^2, \quad (2)$$

where S and T are the dilaton field and the moduli field. Now we assume a nonperturbative superpotential of S and T , $W_{np}(S, T)$, is induced and F -terms of S and T contribute to SUSY breaking. In addition, we assume the vanishing vacuum energy. Then we parameterize F -terms [15]

$$F^S = \sqrt{3}m_{3/2}(S + \bar{S}) \sin \theta e^{-i\alpha_S}, \quad F^T = m_{3/2}(T + \bar{T}) \cos \theta e^{-i\alpha_T}. \quad (3)$$

Within this framework, the soft scalar mass and the A -parameter are obtained

$$m_i^2 = m_{3/2}^2(1 + n_i \cos^2 \theta), \quad (4)$$

$$A_{ijk} = -\sqrt{3}m_{3/2} \sin \theta e^{-i\alpha_S} - m_{3/2} \cos \theta (3 + n_i + n_j + n_k) e^{-i\alpha_T}, \quad (5)$$

where n_i , n_j and n_k are modular weights of fields in the corresponding Yukawa coupling Y_{ijk} . Here we have assumed the corresponding Yukawa coupling Y_{ijk} is T -independent. In addition, the gaugino masses are obtained

$$M_a = \sqrt{3}m_{3/2} \sin \theta e^{-i\alpha_S}. \quad (6)$$

It is obvious that if we require exact degeneracy of squark masses between the first and second families, i.e. $n_{D1} = n_{D2}$ and similar relations for the other squarks, the first 2×2 blocks of the A -matrices for the up and down sector, A_{ij}^u and A_{ij}^d , are degenerate. That is what have been used in Ref. [14]. In such a case we obtain a suppressed contribution to the CP violation ε'/ε .

In the first model we use, we assign different modular weights for the first and second families in order to have asymmetric A matrices. In the dilaton-dominant case with $\tan \theta \gg 1$, we have almost degeneracy of the squark masses. In addition, the renormalization group effects due to the gaugino masses dilute the nondegeneracy. In Ref. [15] it has been shown that the goldstino angle θ is constrained $\cos^2 \theta < 1/3$ for $n_i - n_j = 1$ from the FCNC¹. For example, as our first model with the asymmetric A -matrix, we take the following assignment of the modular weights

$$\begin{aligned} n_{Q_1} &= -1, & n_{Q_2} &= -2, & n_{Q_3} &= -3, \\ n_{D_i} &= n_{U_i} = n_{H_1} = -1, & n_{H_2} &= -3, \end{aligned} \quad (7)$$

where $i = 1, 2, 3$. On the top of that, we restrict our analysis to the region $\cos^2 \theta < 1/3$.

Under this assumption, we have the A -parameter matrix for the down sector,

$$A_{ij}^d = \begin{pmatrix} a_d & a_d & a_d \\ b_d & b_d & b_d \\ c_d & c_d & c_d \end{pmatrix}, \quad (8)$$

where

$$\begin{aligned} a_d &= -\sqrt{3}m_{3/2} \sin \theta, \\ b_d &= m_{3/2}(-\sqrt{3} \sin \theta + e^{-i\alpha'} \cos \theta), \\ c_d &= m_{3/2}(-\sqrt{3} \sin \theta + 2e^{-i\alpha'} \cos \theta). \end{aligned} \quad (9)$$

Here we have rotated the gaugino mass terms into real and rotated the A -terms at the same time and α' denotes $\alpha' \equiv \alpha_T - \alpha_S$. In this case, we have the asymmetry between A_{12}^d and A_{21}^d , but it is limited because of $\cos^2 \theta < 1/3$. Such asymmetry can be enlarged in the h_{ijk} -matrix if we take an asymmetric Yukawa matrix. Thus, we discuss the two cases: one case has a typical symmetric Yukawa matrix and the other case has an example of asymmetric Yukawa matrices. As a typical type of the symmetric and realistic Yukawa matrices, we use the type which are shown explicitly in Ref. [14]. The Yukawa matrices among this type lead to similar results for the CP violation each other. As an example

¹Furthermore, in Ref. [17] it has been shown that a certain type of non-universal A -terms reduce the difference at low energy and even could tune it to vanish.

of asymmetric Yukawa matrices, we take the following form,

$$Y_{ij}^u = y^u \begin{pmatrix} \lambda^8 & \lambda^5 & \lambda^3 \\ \lambda^7 & \lambda^4 & \lambda^2 \\ \lambda^5 & \lambda^2 & 1 \end{pmatrix}, \quad Y_{ij}^d = y^d \begin{pmatrix} \lambda^5 & \lambda^3 & \lambda^3 \\ \lambda^4 & \lambda^2 & \lambda^2 \\ \lambda^2 & 1 & 1 \end{pmatrix}, \quad (10)$$

where $\lambda \sim 0.22$. These correspond to the Yukawa matrices with one $O(\lambda)$ deviation in Ref. [18].

3. We explain the second model which we use. We assume an extra $U(1)$ gauge symmetry² and it is broken by the vacuum expectation value (VEV) of the Higgs field χ . This breaking induces another type of contribution to soft scalar masses, i.e. the D -term contribution, which is proportional to a charge of the broken symmetry. In this case, the soft scalar mass is obtained

$$m_i^2 = m_{3/2}^2(1 + n_i \cos^2 \theta) + q_i m_D^2, \quad (11)$$

where q_i is the charge of the broken $U(1)$ of the matter field, and m_D^2 is the universal part of the D -term contributions. Thus, the soft scalar masses are, in general, non-degenerate for n_i and q_i . However, the soft scalar masses m_i^2 and m_j^2 are degenerate if the following two conditions are satisfied,

$$(a) \quad n_i - n_j = C(q_i - q_j),$$

where C is universal for i and j ,

$$(b) \quad m_{3/2}^2 \cos^2 \theta + m_D^2/C = 0.$$

In this case, we obtain the degenerate soft scalar masses for different n_i and n_j , that is, we can obtain non-degenerate A -matrices keeping degenerate soft scalar masses. Thus, this is a very interesting fine-tuned case in the whole parameter space.

Before calculations of ε'/ε in this model, we give comments on the conditions (a) and (b). We denote here the modular weight and the $U(1)$ charge of χ by n_χ and q_χ and we take the normalization such that $q_\phi = -1$. The VEV of χ induces the following terms in the superpotential,

$$\begin{aligned} W_{Yukawa} &= Y_{ij}^u \theta(q_{H2} + q_{Q_i} + q_{U_j}) (<\chi>/M)^{(q_{H2}+q_{Q_i}+q_{U_j})} Q^i U^j H_2 \\ &+ Y_{ij}^d \theta(q_{H1} + q_{Q_i} + q_{D_j}) (<\chi>/M)^{(q_{H1}+q_{Q_i}+q_{D_j})} Q^i D^j H_1. \end{aligned} \quad (12)$$

The superpotential includes a similar term for the lepton sector. Here the couplings Y_{ij}^u and Y_{ij}^d are naturally of $O(1)$. The suppression factors $(<\chi>/M)^{(q_{H2}+q_{Q_i}+q_{U_j})}$

²This $U(1)$ symmetry may be anomalous [19] or anomaly-free.

and $(\langle \chi \rangle / M)^{(q_{H1}+q_{Qi}+q_{Dj})}$ can lead to realistic hierarchies of the Yukawa matrices [20, 21, 18].

Now we consider the T -duality transformation,

$$T \rightarrow \frac{aT - ib}{icT + d}, \quad (13)$$

where $ad - bd = 1$ and a, \dots, d are integers. We assume that the chiral field transforms

$$\Phi^i \rightarrow (icT + d)^{n_i} \Phi^i. \quad (14)$$

Then we require $G \equiv K + \ln |W|^2$ is duality-invariant. That implies that the superpotential should have the total modular weight $\sum n_i = -3$, that is, [22]³

$$\begin{aligned} (q_{H2} + q_{Qi} + q_{Uj})n_\chi + n_{H2} + n_{Qi} + n_{Uj} &= -3, \\ (q_{H1} + q_{Qi} + q_{Dj})n_\chi + n_{H1} + n_{Qi} + n_{Dj} &= -3, \end{aligned} \quad (15)$$

for non-vanishing couplings. Thus, it is obvious that for Q^i fields the condition (a) is satisfied, i.e.

$$n_{Qi} - n_{Qj} = -n_\chi(q_{Qi} - q_{Qj}), \quad (16)$$

if the (i, k) and (j, k) entries do not vanish for one of k in either the up-sector or down-sector. Similarly we obtain

$$n_{Ui} - n_{Uj} = -n_\chi(q_{Ui} - q_{Uj}), \quad n_{Di} - n_{Dj} = -n_\chi(q_{Di} - q_{Dj}), \quad (17)$$

if (k, i) and (k, j) entries do not vanish for one of k in the up-sector and down-sector, respectively. Thus, a certain type of symmetries can realize the condition (a).

Let us give a comment on the condition (b), too. Now our free parameters are $m_{3/2}$, θ and m_D^2 and these are determined by VEVs. Note that if the condition (a) is satisfied, we have the very special direction in our parameter space, i.e.

$$m_{3/2} \cos^2 \theta = n_\chi m_D^2, \quad (18)$$

which corresponds to the condition (b). Along this direction, we have the degenerate soft scalar masses, $m_{Q(U,D)i} = m_{Q(U,D)j}$ for i, j . For simplicity, we assume that all squark masses are universal for the direction (18), i.e. $m_{Q(U,D)i}^2 = m_0^2$. Now we consider a vicinity around (18) and there are two types of directions to change parameters. One is to violate the degeneracy and the other is to keep the degeneracy. Here let us consider only the

³See also for D -term contributions derived superstring theory e.g. [23].

direction to violate the degeneracy and treat the degree of freedom of the corresponding VEV as a dynamical parameter. Thus, we can write

$$m_i^2 = m_0^2 + q_i \delta m_0^2, \quad (19)$$

around (18), where δm_0^2 corresponds to the dynamical direction to violate the degeneracy. It is obvious that such direction is proportional to q_i for m_i , because of eq.(11) and the linear relation between q_i and n_i .

Now let us consider the one-loop effective potential,

$$\Delta V_1 = \frac{1}{64\pi^2} \text{Str} \mathcal{M}^4 (\ln \mathcal{M}^2 / Q^2 - 3/2). \quad (20)$$

Around (18), we can expand

$$\begin{aligned} \Delta V_1 &= [\Delta V_1]_{m_i^2=m_0^2} + \text{Tr } q_i [d\Delta V_1 / d\delta m_0^2]_{m_i^2=m_0^2} \delta m_0^2 \\ &+ \text{Tr } (q_i)^2 [d^2 \Delta V_1 / d(\delta m_0^2)^2]_{m_i^2=m_0^2} (\delta m_0^2)^2 + \dots \end{aligned} \quad (21)$$

Note that the gaugino masses have no direction corresponding to δm_0^2 . If $\text{Tr } q_i = 0$ and $d^2 \Delta V_1 / d(\delta m_0^2)^2 > 0$, the fine-tuning direction (18) is a locally minimum.

As explained above, we obtain the degenerate soft scalar mass with non-degenerate A -parameters under the conditions (a) and (b). Note that the assignment of the modular weights are related with the assignment of the $U(1)$ charges. As a simple example, we use the following assignment,

$$n_{Q_1} = -4, \quad n_{Q_2} = -3, \quad n_{Q_3} = -1, \quad (22)$$

$$n_{U_1} = -6, \quad n_{U_2} = -3, \quad n_{U_3} = -1, \quad (23)$$

$$n_{D_1} = -3, \quad n_{D_2} = -1, \quad n_{D_3} = -1, \quad (24)$$

$$n_{H_1} = -1, \quad n_{H_2} = -1. \quad (25)$$

This assignment of the modular weights and the corresponding $U(1)$ charge assignment lead to the Yukawa matrices (10) [18]. In realistic string models, we have constraints on the modular weights for the MSSM matter fields [24], and it might be difficult to obtain e.g. the quark field with $n_{U_1} = -6$. However, we use this assignment as a toy supergravity model. If we make our model more complicated e.g. by use of complicated extra symmetries or we concentrate only the first two families, it might be possible to assign more natural values of modular weights.

Namely, our initial conditions in the second model are as follows. We assume $m_i^2 = m_0^2$, and furthermore we assume $m_0^2 = m_{3/2}^2$. We use the Yukawa matrices (10) and the A -matrices (5) with the above assignment of the modular weights.

4. The convenient basis to discuss flavor changing effects in the SUSY loop with gluino exchange is the so-called super-CKM basis [25]. In this basis the relevant quark mass matrix is diagonalized, and the squarks are rotated in the same way. We define $(\delta_{ij})_{LR}$ by normalizing the off diagonal components by average squark mass squared $m_{\tilde{q}}^2$. The relevant contribution, in our models, to CP violation comes from the terms proportional to $(\delta_{12})_{LR}$ and $(\delta_{12})_{RL}$. The mass insertion $(\delta_{12})_{LR}^d$, for instance, is given by

$$(\delta_{12})_{LR}^d = \frac{1}{m_{\tilde{q}}^2} U_{1i}(h^d)_{ij} U_{j2}^T, \quad (26)$$

where U is the matrix diagonalizing the down quark mass matrix. It was shown in Ref. [14] that to obtain a large SUSY contribution to ε it is necessary to enhance the values of $\text{Im}(\delta_{12})_{LR}$ and $\text{Im}(\delta_{12})_{RL}$. The non-degeneracy of A -terms is an interesting example for enhancing these quantities. From eq.(26) we notice that the phases of the all entries of the A -matrix contribute to $(\delta_{12})_{LR}$. It is remarkable that we have $\text{Im}(\delta_{12}^d)_{LR} \neq \text{Im}(\delta_{12}^d)_{RL}$ in the models with asymmetric h -matrices unlike the case of symmetric h -terms. For a wide region of the parameter space we find that $\text{Im}(\delta_{12}^d)_{LR}$ is of order $10^{-4} - 10^{-5}$. That is similar to the case with the symmetric h -terms discussed in Ref. [14]. For these values the CP violation ε , as explained in Ref. [14], is of order of the experimental limit 2.2×10^{-3} . Note that this is a genuine SUSY contribution since we are assuming that $\delta_{CKM} = 0$. Also in this analysis we have assumed the vanishing of the phase of ϕ_μ , which is usually constrained by the experimental limit of the electric dipole moment of the neutron [26, 27].

Now we estimate the $\Delta S = 1$ CP violating parameter ε'/ε in our two models. $(\delta_{12}^d)_{LR}$ and $(\delta_{12}^d)_{RL}$ give the important contributions to the CP violation processes in kaon physics. The relevant part of the effective hamiltonian \mathcal{H}_{eff} for $\Delta S = 1$ CP violation is

$$\mathcal{H}_{\text{eff}} = C_8 O_8 + \tilde{C}_8 \tilde{O}_8, \quad (27)$$

where the Wilson coefficient C_8 and the $\Delta S = 1$ effective Hamiltonian O_8 are given in Ref.[3]. The dependence on $(\delta_{12}^d)_{LR}$ and $(\delta_{12}^d)_{RL}$ appear in C_8 and \tilde{C}_8 ,

$$C_8 = \frac{\alpha_s \pi}{m_{\tilde{q}}^2} \left[(\delta_{12}^d)_{LL} \left(-\frac{1}{3} M_3(x) - 3M_4(x) \right) + (\delta_{12}^d)_{LR} \frac{m_{\tilde{g}}}{m_s} \left(-\frac{1}{3} M_1(x) - 3M_2(x) \right) \right] \quad (28)$$

where $x = m_{\tilde{g}}^2/m_{\tilde{q}}^2$ and \tilde{C}_8 can be obtained from C_8 by exchange $L \leftrightarrow R$ while the matrix element of the operator \tilde{O}_8 is obtained from the matrix element of O_8 multiplying them by (-1) . The functions M_1 , M_2 , M_3 and M_4 can be found in Ref. [3]. The expression for ε' is given by the following formula,

$$\varepsilon' = e^{\frac{\pi}{4}} \frac{\omega}{\sqrt{2}} \xi (\Omega - 1) \quad (29)$$

where $\omega = \text{Re}A_2/\text{Re}A_0$, $\xi = \text{Im}A_0/\text{Re}A_0$ and $\Omega = \text{Im}A_2/(\omega \text{Im}A_0)$. The amplitudes A_I are defined as $\langle \pi\pi(I) | \mathcal{H}_{eff} | K^0 \rangle$ where $I = 0, 2$ is the isospin of the final two-pion state.

Let us discuss the first model with the symmetric and asymmetric Yukawa matrices. We take $\cos\theta = 1/\sqrt{10}$ and $\alpha' \simeq \pi/2$. In both cases, we have $\text{Im}(\delta_{12}^d)_{LR}$ and $\text{Im}(\delta_{12}^d)_{RL}$ of order $10^{-4} - 10^{-5}$ and they are not degenerate. Thus, it is possible to obtain large values of ε'/ε . Indeed for $\text{Im}(\delta_{12}^d)_{LR} \simeq 10^{-5}$ we obtain ε'/ε of order of the experimental results of NA31 and KTeV, $\varepsilon'/\varepsilon = 2.8 \times 10^{-3}$. It is interesting to note that even for $\text{Im}(\delta_{12}^d)_{LR} \sim 10^{-5}$ one can obtain a sizable contribution to ε . As pointed out in Ref. [9], for $\text{Re}(\delta_{12}^d)_{LR} \simeq 10^{-3}$ the ε -requirement $|2\text{Re}(\delta_{12}^d)_{LR}^2 \text{Im}(\delta_{12}^d)_{LR}^2|^{\frac{1}{2}} \simeq 2 \times 10^{-4}$ is satisfied for the region of the parameter space which leads to $\varepsilon'/\varepsilon \sim 10^{-3}$.

Figure 1 shows ε'/ε versus $m_{3/2}$. The solid line and the dotted line correspond to the symmetric and the asymmetric Yukawa matrices, respectively. In the both cases, the SUSY phase has a sizable contribution to ε'/ε . As expected, the asymmetric Yukawa matrix provides with larger values of ε'/ε than the symmetric case. For example, the asymmetric Yukawa matrix leads to the value $\varepsilon'/\varepsilon = 2.8 \times 10^{-3}$ at $m_{3/2} \sim 450$ GeV.

Similarly, we can discuss the second model. Actually, the asymmetry of the A -matrices is large compared with the first model. Thus, the second model predicts much larger values of ε'/ε for the same value of $\cos\theta$. For example, we have $\varepsilon'/\varepsilon = O(10^{-2} \sim 10^{-1})$ for $\cos\theta = 1/\sqrt{10}$ and $m_{3/2} = O(100)$ GeV. Such a case, *i.e.* such magnitude of asymmetry in the h -matrices, is excluded by the KTeV result. We find $\varepsilon'/\varepsilon = O(10^{-3})$ for $\cos\theta = 1/10$. For instance, we have $\varepsilon'/\varepsilon = 3.3 \times 10^{-3}$ for $m_{3/2} = 600$ GeV, $\cos\theta = 1/10$ and $\alpha' \simeq \pi/2$. The behaviour of the $m_{3/2}$ -dependence is similar to the both cases of the first model.

Furthermore, the CP violation ε'/ε can be reduced if any elements in the Yukawa matrices (10) include suppression factors such that the asymmetry of the h -matrices is reduced. For example, here we replace the (3,1) element of the down Yukawa matrix as $\lambda^2 \rightarrow x\lambda^2$. For $x = 0.1$ we have $\varepsilon'/\varepsilon = O(10^{-3})$ in the case with $\cos\theta = 1/\sqrt{10}$ and

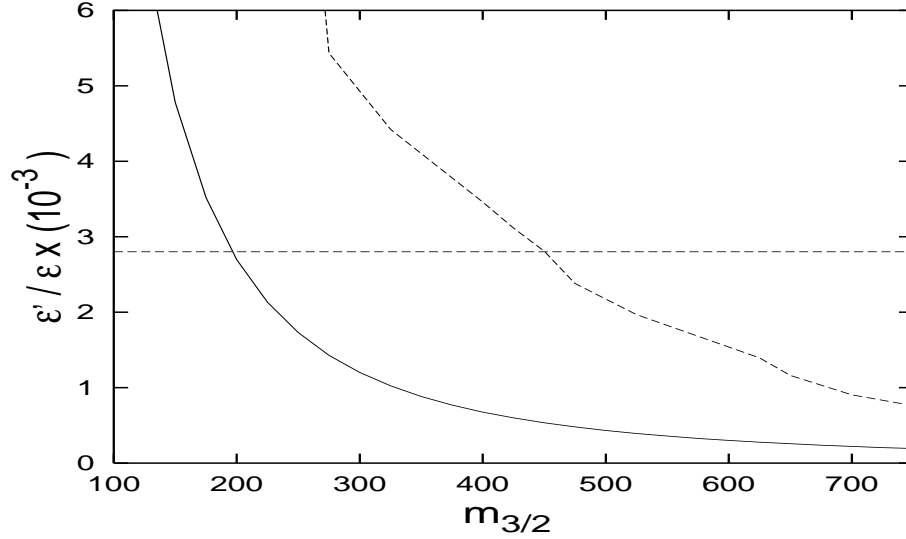


FIGURE 1. The value of ε'/ε as a function of the gravitino mass

$m_{3/2} = O(100)$ GeV. For instance, we find $\varepsilon'/\varepsilon = 2.5 \times 10^{-3}$ for $x = 0.1$, $\cos \theta = 1/\sqrt{10}$ and $m_{3/2} = 300$ GeV.

5. We have studied the CP violation ε'/ε in the models with asymmetric A -matrices as well as asymmetric h -matrices. We have shown that a certain type of the asymmetry enhances ε'/ε and it can be of order of the KTeV result, $\varepsilon'/\varepsilon \sim 10^{-3}$. A large magnitude of the asymmetry leads to too large CP violation ε'/ε . Thus, we have a constraint on the large asymmetry of h -matrices from the experimental value $\varepsilon'/\varepsilon \sim 2.8 \times 10^{-3}$. Other CP violating aspects and FCNC processes would also give further constraints.

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